



## THE STATUS OF ELECTROWEAK GAUGE THEORY

Dallas C. Kennedy

*Fermilab, P.O. Box 500 MS-106  
Batavia, Illinois 60510-0500 USA*

### ABSTRACT

The present status of electroweak gauge interactions, measured in flavor-conserving fermionic processes, is discussed. The Standard Model gauge structure  $SU(2)_L \times U(1)_Y$  has now been tested, with the advent of the LEP experiments, to a few tenths of a percent, including the effect of radiative corrections. A special feature of electroweak gauge theory, the non-decoupling of heavy virtual states with weak quantum numbers, is now highly constrained by the data, placing serious restrictions on the types and extent of non-Standard sectors heavier than the  $Z$  boson.

*Invited talk presented at the 21<sup>st</sup>  
Coral Gables/Global Foundation Conference:  
Unified Symmetry in the Small and in the Large,  
January 25 - 27, 1993*

### 1. ELECTROWEAK STANDARD MODEL [1]

The Standard Model of electroweak gauge interactions, developed by Glashow, Salam, and Weinberg in the 1960's and for which they received the Nobel prize [2], has now withstood its most stringent tests yet without failing. Based on the non-Abelian gauge group  $SU(2)_L \times U(1)_Y$  — weak left isospin mated to weak hypercharge — this theory incorporates the highly successful quantum electrodynamics [3], with the gauge subgroup of electric charge  $U(1)_Q$ , is unitary and renormalizable, and reproduces all known weak neutral- and charged-current phenomenology to a fraction of a percent [4]. Here, I report on the status of this theory, concentrating on the flavor-conserving, four-fermion, CP-conserving gauge exchange sector.<sup>1</sup>

The gauge sector of the electroweak Standard Model is characterized by three (or four, depending on assumptions about the symmetry breaking) parameters. The theory, based on a product group of Lie algebras, has two independent gauge couplings, the  $SU(2)_L$   $g$  and the  $U(1)_Y$   $g'$ . These are usually expressed in terms of the proton's electric charge  $e$ :  $1/e^2 = 1/g^2 + 1/g'^2 = 1/4\pi\alpha$ , and the weak mixing angle  $\tan\theta_W = g'/g$ . The third parameter sets the mass scale for the theory, the Higgs vacuum expectation value  $v$ . The Higgs sector is a collection of scalar fields (at least

<sup>1</sup>I therefore ignore flavor-changing neutral currents and CP violation, subjects treated elsewhere in these proceedings.



one) that acquire non-zero vacuum condensates and thereby spontaneously break the full gauge symmetry down to  $U(1)_Q$  [5]. This quantity is usually quoted in terms of the Fermi weak constant  $G_F = 1/\sqrt{2}v^2$ , where  $v = 246$  GeV. A possible fourth parameter is the  $\rho$  of Veltman, which characterises the vacuum structure. In the Minimal Vacuum Standard Model (MVSM), the Higgs sector contains only weak doublets, and  $\rho = 1$  automatically. If we relax this requirement and allow non-doublets in the vacuum, then  $\rho$  is arbitrary and must be experimentally measured. I call this case the Extended Vacuum Standard Model (EVSM). These three (or four) numbers are sufficient to specify the electroweak gauge interactions completely at tree level.

The tree approximation is adequate to predict electroweak interactions to about the percent level. The theory contains three basic types of gauge interaction: the parity-conserving, long-range electromagnetic neutral and massless photon; and the massive, short-range weak neutral- and charged-current  $Z$  and  $W$  bosons. The charged-current  $W$  couples to ordinary fermions solely by left-handed currents and is responsible for the familiar low-energy weak beta decay. The neutral-current  $Z$  coupling to ordinary fermions is also parity-violating, with strength controlled by the weak mixing  $\sin\theta_W$ . A paradigmatic low-energy  $Z$  exchange process is neutral-current neutrino scattering. The electromagnetic and weak couplings of the fermions have been measured to nearly exact universality, by experiments ranging from very low energies to  $W$  and  $Z$  production at the LEP, SLC, and Tevatron colliders. This is now the thoroughly investigated “low-energy” world of four-fermion processes, beyond which lies the unexplored realm of the very high-energy colliders, HERA, LEP200, LHC, SSC, and possibly others. What might we find there?

Two minimally necessary pieces of the Standard Model still undiscovered are the Higgs boson(s) and the top quark. Recall that a Higgs sector is needed to break the electroweak gauge symmetry. The minimal Standard Model contains one complex doublet, with four degrees of freedom; three are would-be Goldstone bosons disguised as the longitudinal  $W^\pm$  and  $Z^0$  degrees of freedom, leaving one physical Higgs boson, a neutral scalar isoscalar. At least this one state is required to preserve the unitarity and renormalisability of the theory [4]. The Higgs mass is proportional to the vacuum expectation value, the proportionality constant being an unknown Higgs four-point self-coupling. If the Higgs is light to moderately heavy (less than about 600 GeV), it is said to be *weakly coupled*, *physical*, or *linear*, with a decay width modest compared to its mass. If the Higgs is heavy (more than 600 GeV), then the self-coupling is large, and the Higgs sector is called *strongly coupled*, *unphysical*, or *nonlinear*, because the decay widths are typically so large that one-particle states lose distinct meaning. The Higgs boson in that case becomes analogous to the elusive hadronic  $\sigma$  resonance. Lattice simulations indicate that the Higgs mass probably cannot be larger than about 800 GeV; at that point, the Higgs sector is highly nonlinear and saturated [6].

The top quark is the partner of the bottom quark, in a weak left isospin doublet, where  $I_3^L(b) = -1/2$  and  $I_3^L(t) = +1/2$ . The existence of the top is unavoidable for symmetry reasons: the bottom isospin has now been measured [7], and its doublet companion is required for renormalisability and anomaly cancellation. The top quark is too heavy to be produced as yet, but, as we see later, its virtual radiative effects can be used to limit its mass to less than about 200 GeV.<sup>2</sup> The masses of ordinary fermions are proportional to the Higgs vacuum expectation  $v$ , through their Yukawa couplings to the Higgs sector. The top must have a large Yukawa coupling. Thus the properties of the top, especially its decay modes and partial widths, may provide non-trivial information about the Higgs sector [9].

## 2. RADIATIVE CORRECTIONS TO GAUGE INTERACTIONS [1]

The general precision of electroweak data is now at the level of a few tenths of a percent. At such detail, meaningful comparisons between theory and experiment require the inclusion of radiative

<sup>2</sup>The CDF and DØ collaborations at the Fermilab Tevatron report an official production lower bound of 91 GeV [8] for the top quark mass and a current unofficial lower limit of 108 GeV (March 1993).

corrections, usually only through one loop. The radiative corrections to electroweak gauge exchange fall into two broad classes [10, 11].

The first class contains the corrections that have the same Lorentz and gauge symmetry properties as the tree- or Born-level interactions. They have been variously called *oblique*, *universal*, or *Born-like*. Such corrections contain ultraviolet divergences, but, because they reproduce the same form as the tree-level interactions, they simply renormalise the tree-level parameters in a momentum-dependent way, and the associated divergences are absorbed into the redefinition of parameters. Following the system of Kennedy and Lynn, I denote the effective running or momentum-dependent parameters with a  $*$  subscript:  $e_*^2(q^2)$  replaces  $e^2$ ,  $s_*^2(q^2)$  replaces  $\sin^2 \theta_W$ ,  $G_{F*}(q^2)$  replaces  $G_F$ , and  $\rho_*(q^2)$  replaces  $\rho$ , with  $q^2$  the invariant square-momentum transfer of the given process. (Even if  $\rho \equiv 1$  at tree level,  $\rho_*(q^2)$  is not unity in general; see below.) The universal corrections contain all the gauge boson self-energy corrections, plus selected parts of the vertex and box corrections necessary to maintain gauge invariance in a non-Abelian theory. Gauge invariance also guarantees that the Born-like corrections are truly universal, as outlined by Degrassi and Sirlin [12], even in non-fermionic processes. Such corrections should be thought of as intrinsic properties of the gauge bosons, independent of the identities of the external particles.

The second class are the remaining vertex, box, and bremsstrahlung (radiation) corrections that do not reproduce the Lorentz and gauge symmetry properties of the tree-level interactions. These go under the names of *direct*, *non-universal*, or *non-Born-like* corrections. Because they cannot renormalise a tree-level parameter, they are necessarily finite. They also depend on the specific quantum numbers (mass, spin, charge, isospin) of the external particles.

With one exception in section 5, I concentrate exclusively on the universal corrections. Because they are the same in all gauge interactions, we can pool together the data from many different processes to test this sector of the electroweak theory, taking into account only the momentum transfer of each process. The universal corrections come in two types. The first type consists of corrections that renormalise the gauge couplings  $e_*^2(q^2)$  and  $s_*^2(q^2)$ . These are just generalisations of the photon vacuum polarisation in quantum electrodynamics. In particular, they respect the decoupling theorem of Appelquist, Carazzone, and others [13]: at low energies,  $|q^2| \ll \text{heavy mass-squares } M^2$ , the effect of heavy particles in the loops decouples as  $\mathcal{O}(q^2/M^2)$  or faster. In section 3, I represent the running gauge couplings in the more conventional  $\overline{MS}$  scheme, which is convenient because it absorbs the leading (logarithmic) momentum dependence of the full running functions  $e_*^2(q^2)$  and  $s_*^2(q^2)$ . The second type of universal corrections renormalise the running symmetry-breaking parameters  $G_{F*}(q^2)$  and  $\rho_*(q^2)$ . The decoupling theorem does not apply to a *broken* gauge theory; the *non-decoupling* effects of heavy virtual particles appear in this second type of correction. Such effects generally vary with quadratically, logarithmically, or as constants with the heavy particle masses [14, 15, 16, 10, 11].<sup>3</sup>

These non-decoupled loop effects can place constraints on physics heavier than the  $Z$  boson, if we make a few broad assumptions. The reason that the gauge symmetry must be broken is that the non-decoupled effects arise in the longitudinal gauge degrees of freedom, which are Goldstone bosons in disguise. For example, the gauge boson self-energy in an unbroken theory is proportional to  $q^2$  and thus vanishes at  $q^2 = 0$ , corresponding to zero mass for the gauge boson. In a broken theory, the self-energy does not vanish at  $q^2 = 0$ , nor does it need to, since the gauge mass is already non-zero to begin with. Such components of the self-energy, in effect, renormalise the longitudinal gauge modes = Goldstones and the gauge masses. They are really Goldstone scalar self-energies.

<sup>3</sup>The original electroweak renormalisation scheme in common use was the *on-shell scheme*, introduced by Sirlin and Marciano in 1980 [17], with the mixing  $\sin^2 \theta_W \equiv 1 - M_W^2/M_Z^2$  to all orders of perturbation theory. While this scheme is convenient for low-energy processes, it is quite difficult to use when pooling together processes of widely differing energies, because it mixes running gauge coupling and non-decoupled heavy physics corrections together in a single parameter,  $\sin^2 \theta_W$ , by defining it in terms of the gauge boson masses. Consequently, Sirlin introduced the  $\overline{MS}$  scheme into electroweak radiative corrections in 1989 [18], where it dovetails nicely with the universal  $*$  gauge coupling functions.

Thus, to see non-decoupled effects in a gauge theory, the gauge symmetry must be broken, so that the gauge bosons can mix with the Goldstones; and the heavy particles in the loops must acquire their masses from the *same* Higgs sector that breaks the gauge symmetry.

But there is a third requirement. When the gauge boson masses are renormalised, the non-decoupled effects can still be cancelled by a subtraction and be absorbed into an arbitrary parameter. The only way to prevent this without fine-tuning these parameters is to have an approximately good global symmetry that holds at tree level but is violated by loop effects. Then observed deviations from the exact symmetry can be attributed to loop effects and separated out [19]. For the electroweak theory, the relevant global symmetry was discovered by Veltman [14], Appelquist and Bernard, Longhitano, and others [20], and occurs in the Higgs/Goldstone sector, the global weak chiral group  $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ , left times right times baryon minus lepton number. This group is generation-independent. The subgroup  $SU(2)_L \times U(1)_Y$ , where the hypercharge  $Y/2 = I_3^R + (B - L)/2$ , is the gauge group; the electric charge is the unbroken, parity-conserving  $Q = I_3^L + Y/2 = I_3^L + I_3^R + (B - L)/2$ . By separating out the non-decoupled effects in gauge interactions, we uncover some picture of the underlying Higgs sector through the longitudinal gauge modes.

A well-known example of a non-decoupled radiative effect made possible by the breaking of a global symmetry is the appearance of flavor-changing neutral currents (FCNCs) in the Standard Model [21]. The relevant group is the Glashow-Iliopoulos-Maiani (GIM) family symmetry, which rotates all the up-type quarks ( $u, c, t$ ) into one another; and similarly all the down-type quarks ( $d, s, b$ ). The electroweak gauge sector respects this symmetry, which forbids flavor-changing neutral currents. However, the Higgs sector clearly breaks the symmetry, since the up and charm quark masses are different, and so on. Thus at loop level, it becomes possible to have flavor-changing neutral currents, depending on the mass differences among the up-type and down-type quarks, respectively. The global GIM symmetry forbids this effect at tree level, while at loop level is unable to prevent it, but nevertheless keeps the flavor-changing neutral current loops finite.

We know little about the Higgs sector, except that the electroweak gauge symmetry must be broken by something that leaves the Lorentz and electric charge symmetries intact. We do know one more important fact. (I briefly revert to tree-level notation.) The  $Z$  boson is a mixture of two unphysical neutral states, the neutral gauge modes  $W^0$  of  $SU(2)_L$  and  $B$  of  $U(1)_Y$ . That is,  $Z = \cos \theta_W \cdot W^0 + \sin \theta_W \cdot B$ . Now,  $M_W = \rho M_Z \cos \theta_W$ , and  $\rho$  is very close to unity, implying that  $M_{W^\pm} \simeq M_{Z^\pm}$ . As noted by Veltman [14], the only natural way for this to hold is if the vector subgroup  $SU(2)_V$  of  $SU(2)_L \times SU(2)_R$  remains unbroken, where  $I_V = I_L + I_R$ . The overall weak chiral group must be broken to accommodate the Higgs vacuum expectation value, but the vector or *custodial* subgroup need not be. The easiest way to guarantee this requirement, in turn, is for all the Higgs scalars to be weak doublets. Thus, an alternative and more general definition of the MVSM is that *the weak custodial subgroup be unbroken by the vacuum at tree level*. In this case, the slight deviation of the measured  $\rho$  from unity is attributed to custodial-breaking loop effects. The EVSM case, then, assumes that *the custodial subgroup is also broken at tree level*, allowing for an arbitrary  $\rho$  parameter. The deviation of  $\rho$  from unity is one possible non-decoupled loop effect. Since the  $W$  and  $Z$  bosons have non-zero masses, they can have non-trivial on-shell wavefunction renormalisations, making possible two other non-decoupled effects. In all, the MVSM case thus has three distinct generic parameters for non-decoupled loop effects. In the EVSM case, one of these loop effects disappears and reappears as the tree-level  $\rho$  parameter, whose value is *a priori* arbitrary [11, 19].

The functional information concerning all possible non-decoupled universal radiative corrections is contained in the three finite combinations of gauge self-energies derived by Kennedy and Lynn [11], generalising the earlier work of Lynn, Peskin, and Stuart [10]:

$$\begin{aligned}\Delta_\rho(q^2) &= \Pi_\pm(q^2) - \Pi_{33}(q^2), \\ \Delta_3(q^2) &= \Pi_{33}(0) + \Pi_{3Q}(q^2) - \Pi_{33}(q^2), \\ \Delta_\pm(q^2) &= \Pi_\pm(0) + \Pi_{3Q}(q^2) - \Pi_\pm(q^2),\end{aligned}\tag{1}$$

where the subscripts refer to left isospin (3, +, -) and charge ( $Q$ ) currents. For theoretical convenience and experimental reasons, it proves useful to introduce equivalent dimensionless parameters, such as those proposed by Peskin and Takeuchi [22, 19]:

$$\begin{aligned}\alpha T &= 4\sqrt{2}G_F\Delta\rho(0), \\ S &= -16\pi\Delta_3(M_Z^2)/M_Z^2, \\ S+U &= -16\pi\Delta_\pm(M_W^2)/M_W^2,\end{aligned}\tag{2}$$

with an indefinite number of equivalents [23]. The specific momentum points are chosen because of the theoretical and experimental simplifications that occur at the gauge boson poles and in the low-energy weak neutral current, although the parameters could also be expressed as derivatives of self-energies or as coefficients of a momentum expansion. The parameter  $T$  simply gives the deviation of  $\rho$  from unity, while  $S$  and  $S+U$  are the  $Z$  and  $W$  wavefunction renormalisations, respectively. Their symmetry properties under the global group are straightforward [19]. The overall  $SU(2)_L \times SU(2)_R$  symmetry is broken by the static Higgs vacuum at  $q^2 = 0$  (and fixed by the value of  $G_F$ ), but this leaves open dynamical or  $q^2$ -dependent effects that can break the overall global group in loops; this effect is measured by the parameter  $S$ . In the MVSM case, the custodial subgroup  $SU(2)_V$  is unbroken at tree level; at the loop level, static and dynamical effects breaking  $SU(2)_V$  are possible, given by  $T$  and  $U$ , respectively. In the EVSM case (if we assume Higgs non-doublets),  $T$  disappears into the tree-level  $\rho$ , according to the prescription:  $(1 - \alpha T)^{-1} \rightarrow \rho$ .

The  $T$  parameter effects have been familiar for some time, as  $\rho$  effects. Thus, the top-bottom mass splitting (which breaks the  $SU(2)_V$  symmetry) contributes to  $T$  as:

$$\alpha T \simeq 3G_F m_t^2 / 8\sqrt{2}\pi^2,\tag{3}$$

for large top mass, while the massive Higgs contributes as:

$$\alpha T \simeq -(3G_F(M_Z^2 - M_W^2)/8\sqrt{2}\pi^2)\ln(m_H^2/M_Z^2),\tag{4}$$

depending on the custodial-breaking  $Z - W$  splitting [14]. The  $S$  and  $U$  effects, on the other hand, have not been widely recognised until recently, because they require  $Z$  and  $W$  pole measurements [23]. The leading dependences on the top and Higgs masses are:

$$\begin{aligned}S &\simeq -(1/6\pi)\ln(m_t^2/M_Z^2) + (1/12\pi)\ln(m_H^2/M_Z^2), \\ U &\simeq +(1/2\pi)\ln(m_t^2/M_Z^2),\end{aligned}\tag{5}$$

where, for all three parameters, I take  $m_t = m_H = M_Z$  as my reference point for the minimal Standard Model.

### 3. PRECISION ELECTROWEAK MEASUREMENTS<sup>4</sup>

It is possible to develop the above theoretical structure and compute heavy physics contributions to electroweak radiative corrections without any experimental input. But to carry out precision tests of electroweak gauge interactions requires precision data. Since the early 1980's, such data have accumulated to test the electroweak theory at the radiative level. The pre-1988 data include neutrino-nucleus and neutrino-electron scattering, electron-positron annihilation below the  $Z$  pole, and polarized charged lepton-nucleus scattering. (This data collection was summarized by Amaldi *et al.* [24].) Since the beginning of 1988, new categories of data have appeared: the  $Z$  mass, width, and asymmetry measurements of LEP (CERN) and SLC (SLAC); the  $W$  mass measurements of CDF (Tevatron/Fermilab) and UA2 (Sp $\bar{p}$ S/ CERN); and the atomic parity violation measurement

<sup>4</sup>The data and fits shown here are current through September 1992.

in cesium at NIST/U. of Colorado (Boulder) [25]. At the same time, better neutrino scattering data have appeared from CHARM-II (CERN) and CCFR (Fermilab). These are all four-fermion processes, and we assume, with an exception in section 5, only universal radiative corrections, so that all these data can be combined into a single fit. Table 1 lists the common pool of available electroweak gauge interaction data. The complete minimal Standard Model corrections (bremsstrahlung where necessary, vertices, and boxes) are also included.

Data	Process	Experiment
Pre-1988	$(\nu, \bar{\nu})N$ NC	CCFR, FMM (FNAL)
		CHARM, CDHS (CERN)
	$(\nu, \bar{\nu})(n, p)$ NC	FNAL, CERN, BNL
	$\nu N \rightarrow \nu \pi^0 N$	FNAL, CERN, Serpukhov
	$(\nu, \bar{\nu})e$	BNL, LANL, SRR, FNAL, CERN
	$I^\pm(\text{pol})N$	SLAC, CERN
	$e^+e^- < Z$	SLAC, DESY
Post-1988 $\nu$	$\nu_\mu e / \bar{\nu}_\mu e$	CHARM-II (CERN)
	$\nu N$ NC	CCFR (FNAL)
Atomic Parity	$6S/7S - 6P$ mixing $^{55}\text{Cs}$ [133]	Colorado/NIST
$W$ Mass	$p\bar{p} \rightarrow W$	CDF (FNAL), UA2 (CERN)
$Z$ Resonance	$e^+e^- \rightarrow Z$	LEP (CERN)
	$Z$ mass, widths, FB asymm,	"
	$r(\text{pol})$ asymm	"
	$e^+e^-(\text{pol})$ asymm†	SLC (SLAC)

Table 1. Electroweak gauge data from four-fermion processes (September 1992). † Preliminary.

There are five basic observables that these data measure [23]. The first two are the weak gauge boson masses. Let  $M_{ZO}$  and  $M_{WO}$  be these masses if there were no heavy physics. Then the actual masses are:

$$\begin{aligned}
 M_Z^2 &= M_{ZO}^2 \left[ \frac{1 - \alpha T}{1 - 4\sqrt{2}G_F M_{ZO}^2 S/16\pi} \right], \\
 M_W^2 &= M_{WO}^2 \left[ \frac{1}{1 - 4\sqrt{2}M_{WO}^2(S+U)/16\pi} \right].
 \end{aligned} \tag{6}$$

The next two are the weak gauge boson widths, evaluated at their respective poles and renormalized by their respective pole residues:

$$\begin{aligned}
 \Gamma_Z &= G_F M_Z^3 \gamma_Z / (1 - \alpha T), \\
 \Gamma_W &= G_F M_W^3 \gamma_W,
 \end{aligned} \tag{7}$$

where  $\gamma_Z$  is the canonical function of  $s_z^2(Z)$  without any explicit heavy physics. Finally, we have the low-energy matrix element for the weak neutral current exchange between fermion lines of isospin and charge  $I_L^2, Q$  and  $I_L^{2'}, Q'$ , respectively:

$$\mathcal{M}_{NC}(q^2 \rightarrow 0) = \frac{4\sqrt{2}G_F}{1 - \alpha T} [I_L^2 - s_z^2(0)Q] [I_L^{2'} - s_z^2(0)Q'], \tag{8}$$

where the reciprocal factor of  $1 - \alpha T$  is the familiar  $\rho$  parameter. Let us pause to note some interesting points. The low-energy neutral-current exchange (8) depends universally on  $s_z^2(0)$  and

$T$  only. Hence, low-energy measurements alone tell us about the  $\rho$  parameter, but no other heavy physics. Furthermore,  $T$  appears in neutral-current processes only. The relation between  $s_w^2(0)$  and the canonical  $s_w^2(Z)$  is a gauge coupling running and thus does *not* depend on heavy physics. However, if we then express  $s_w^2(Z)$  in terms of  $M_Z$ , this relation, because it explicitly brings in electroweak symmetry breaking, *does* depend on heavy physics. With  $M_Z$  used as an input,  $s_w^2(Z)$  is no longer a free parameter, but is determined by  $M_Z$ ,  $S$ , and  $T$ , and is in fact overdetermined by the data, a feature that gives some check on the self-consistency of the  $SU(2)_L \times U(1)_Y$  gauge structure. Also, note that the parameter  $U$  appears only in the  $W$  mass. The low-energy neutral-current interaction in atomic parity violation has a unique feature. The underlying electroweak quantity measured is the so-called *weak charge* of the nucleus, obtained from the electron axial-vector coupling times the nuclear vector coupling and depending on  $s_w^2(0)$ . If  $s_w^2(0)$  is re-expressed in terms of  $s_w^2(Z)$  and in turn in terms of  $M_Z$ ,  $S$ , and  $T$ , then  $\mathcal{M}_{NC}(0)$  contains new, implicit dependence on  $S$  and  $T$ , besides the explicit  $\rho$  factor. In heavy hydrogenic atoms such as cesium, where the proton/neutron ratio is about  $2/3$  (because of Coulomb repulsion in the nucleus) and with  $s_w^2(0) \simeq 1/4$ , the  $T$  dependence in the overall parity-violating interaction between the nucleus and the single valence electron cancels almost exactly. The cesium experiments are thus sensitive to  $S$ , but not  $T$  [23]. The now-standard precisely-known tree-level inputs are:  $M_Z$ , determined from the peak of the  $Z$  resonance;  $G_F$ , determined from muon beta decay; and  $\alpha$ , determined from low-energy electrodynamics.

The results for the weak mixing (in the  $\overline{MS}$  scheme) and the non-decoupled parameters are [26]:

$$\begin{aligned}\sin^2 \hat{\theta}_W(M_Z^2) &= 0.2313 \pm 0.0010, \\ S &= -0.8 \pm 0.5, \quad T = -0.2 \pm 0.4, \\ U &= +0.2 \pm 0.9,\end{aligned}\tag{9}$$

in the MVSM case. The 90 (95)% C.L. upper bounds on these quantities are:  $S < -0.1$  (0.0),  $T < 0.4$  (0.5), and  $U < 1.4$  (1.7). The  $Z$  mass is  $M_Z = 91.187 \pm 0.007$  GeV, while the muon beta decay Fermi constant is  $G_\mu = (1.16637 \pm 0.00002) \times 10^{-5} \text{GeV}^{-2}$ . The value of  $s_w^2(Z)$  is assumed, calculated from the low-energy  $\alpha$  with vacuum polarization applied, including the hadronic component computed via dispersion relation from the  $e^+e^- \rightarrow \text{hadrons}$  data:  $1/\hat{\alpha}(M_Z^2) = 127.8 \pm 0.1$ , with hadronic uncertainty [18]. If we substitute in the explicit dependences of the theory on the top quark and Higgs boson masses, we obtain a minimal Standard Model fit:

$$\begin{aligned}\sin^2 \hat{\theta}_W(M_Z^2) &= 0.2325 \pm 0.0007, \\ m_t &< 193 \text{ (201) GeV, 90 (95)\% C.L.} \\ &= 158 \pm 27 \text{ GeV,}\end{aligned}\tag{10}$$

including the  $\mathcal{O}(\alpha_s)$  one-gluon exchange correction to the parameter  $T$  [27]. The Higgs mass dependence, which is only logarithmic, is so weak that no useful limits on  $m_H$  can be inferred. The uncertainties in the top quark mass and in the weak mixing include allowing the Higgs mass to range from 50 to 1000 GeV. Figure 1 shows the  $(S, T)$  plane with the 90% C.L. limit region shaded (with  $U$  set to zero by hand to reduce the fit to two dimensions). The calibrated curve in the upper half of the plane gives the Standard Model dependence on the top quark mass. Note in particular that the accuracies of the  $Z$  width and asymmetry measurements are now so good that they exert the most control on the final result of the fit; the atomic parity violation measurement, which is sensitive to  $S$  alone, has such a large uncertainty that it barely affects the outcome. Figure 2 shows the top quark mass 90% C.L. limit alone from various data constraints. Note that the stringent limit on the top quark mass arises from the explicit appearance of  $T$  in the fit; in turn, this requires the minimal vacuum assumption, so that  $\rho = 1$  at tree level. Note also, as we should expect, that atomic parity violation, lacking any  $T$  dependence, does not contribute to the top quark mass bound.

It is also possible to use the precise value of  $\sin^2 \theta_W$  to test models of grand unification [28, 29], since grand unified theories (GUTs) predict its value, rather than take it as an arbitrary input.

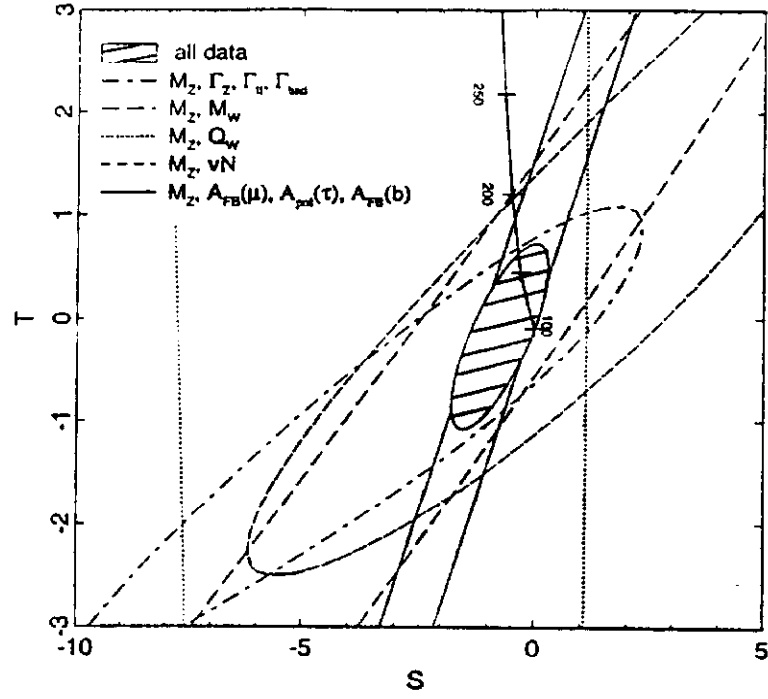


Figure 1. Combined electroweak data fit for  $(S, T)$  with  $U \equiv 0$ . 90% C.L. region is shaded [23, 26].

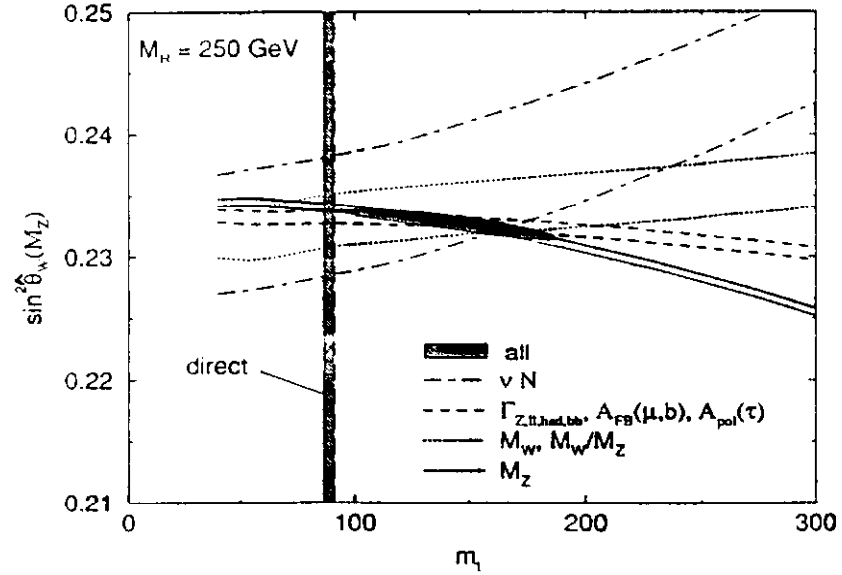


Figure 2. Combined electroweak data fit constraining top quark mass  $m_t$  (in GeV) in the minimal Standard Model,  $m_H = 250$  GeV [28, 26].



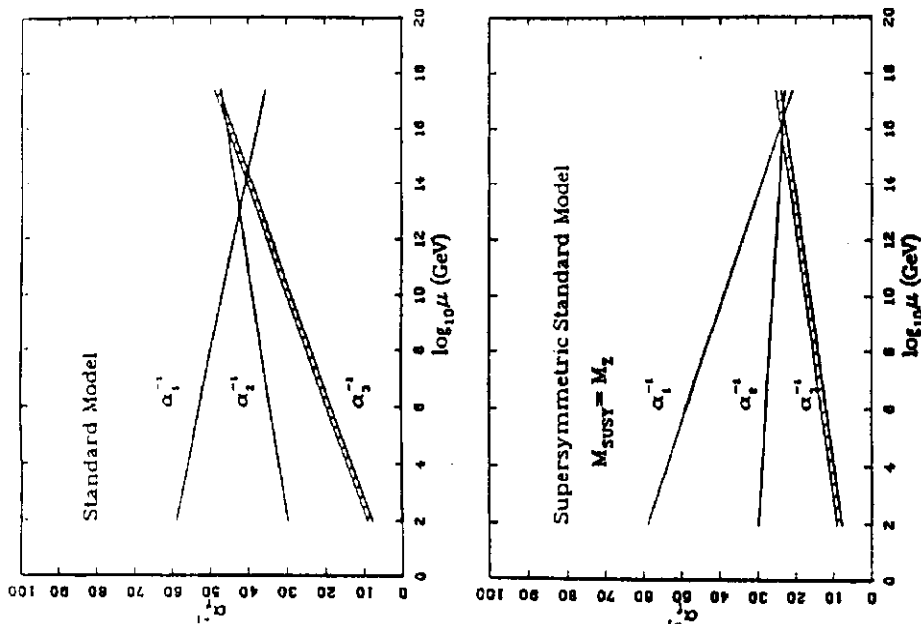


Figure 3. Running of three Standard Model gauge couplings to unification: (a) minimal Standard Model embedded in minimal  $SU(5)$  GUT; (b) minimal supersymmetric Standard Model embedded in supersymmetric  $SU(5)$  GUT [28, 29].

Figure 3 shows the running together of the three gauge couplings,  $SU(3)_C$ ,  $SU(2)_L$ , and  $U(1)_Y$ , under two sets of assumptions. The upper graph assumes no new states beyond the minimal Standard Model until the unification scale  $M_X$ , the so-called “Great Desert” scenario; the Standard Model then unifies into the minimal  $SU(5)$  GUT at  $M_X$ . But this model fails to unify the three couplings at the same scale. (It is also already known to give too short a proton lifetime.) On the other hand, as we see in the lower graph, if we add to the minimal Standard Model its complement of supersymmetric partners at some scale between 100 and 1000 GeV, the unification of the three couplings proceeds without difficulty. (This includes uncertainties due to threshold effects and non-renormalisable operators.) The unification scale,  $M_X \simeq 2 \times 10^{16}$  GeV, is consistent with current bounds on the proton lifetime, although some of the proton decay modes predicted by supersymmetric  $SU(5)$  should be visible in the next generation of detectors [31]. If we turn grand unification around and *assume* a group to start with, then the electroweak couplings alone,  $e^2(Z)$  and  $g^2(Z)$ , are sufficient to determine the unification scale and predict, rather than take as input, a value of the strong  $SU(3)_C$  coupling. This approach has the nice feature of avoiding strong interaction uncertainties in the GUT inputs [30]. The strong coupling at the  $Z$  pole now stands at  $\hat{\alpha}_s(M_Z^2) = 0.12 \pm 0.01$ , with a conservative uncertainty, inferred from a combination of  $Z$  hadronic decay branching ratios and  $Z \rightarrow \text{jets}$  event topologies.

#### 4. IMPLICATIONS FOR HEAVY PHYSICS

If the MVSM assumption is correct (see section 5), the top quark mass limit is good news for the Tevatron, since top is very likely to be visible at the CDF and D0 detectors [9, 32]. (If the top quark is *not* seen below 200 GeV, then the vacuum is probably extended.) More generally, the fact that  $T$  is close to zero implies that, unless we appeal to special cancellations, the custodial isospin-breaking

corrections are small, with mass splittings about the same as the limit on the top quark mass. The limit on  $S$  is more interesting, inasmuch as most conceivable kinds of non-Standard physics contribute *positively* to  $S$ , while the favored range for  $S$  is negative. A degenerate Dirac fermion left isodoublet, for example, contributes  $+1/6\pi$  to  $S$ , so that a full degenerate fourth generation would contribute  $+4/6\pi = +0.21$  [11, 19]. If we saturate the  $S$  bound with nothing else, even one degenerate new generation is ruled out at two standard deviations. Negative values of  $S$  are possible with heavy Majorana neutrinos [33] and certain scalar mass hierarchies [34].

The major rivals for non-Standard physics are technicolor [35] (dynamical electroweak symmetry breaking) and supersymmetry [36], both introduced primarily to cure the radiative instabilities of the Higgs sector (the gauge hierarchy problem). Technicolor cures the instabilities by replacing elementary scalars with composites of technifermions bound by a new strong force. Supersymmetry retains elementary scalars, but relates them through supersymmetry to fermionic superpartners (Higgsinos) whose masses are protected from radiative instabilities by chiral symmetry. The main bound on technicolor theories derives from  $S$ , as this seems to be computable, in some approximation, in conventional technicolor theories. By conventional, I mean strong sectors with vector-like, unbroken, and confining gauge groups. This category includes both scaled-up QCD-like theories and the *walking* technicolor (WTC) theories [37], whose running couplings are nearly constant to very high energies (50-100 TeV). The contributions to  $S$  in the QCD-like theories arise from three sources: vector (technirho) resonances, axial-vector (techni- $A_1$ ) resonances, and pseudoGoldstone bosons (technipions), with the first contribution being the main one. The last set was calculated by Golden and Randall [38], while the full calculation was done by Peskin and Takeuchi [22], in the large- $N$   $SU(N_{TC})$  limit, using the known sum rule properties of QCD as an analog computer for  $N_{TC} = 3$ :

$$S \simeq (0.10)N_{TC}N_{TD} + 0.13, \quad (11)$$

where  $N_{TC}$  is the number of technicolors,  $N_{TD}$  the number of technifermion electroweak doublets, and the final term reflects taking the effective Higgs mass to a TeV. This result should be compared with the naive fermion loop result,  $N_{TC}N_{TD}/6\pi = (0.05)N_{TC}N_{TD}$ . Clearly, the QCD-like theories are disfavored by the electroweak gauge measurements. But these theories were already known to produce problems with flavor-changing neutral currents which WTC was introduced to cure. The parameter  $S$  is harder to compute in these theories, because not much is known about their spectra and they lack the analogy to QCD. Nonetheless, a number of workers have attempted to estimate  $S$  and have arrived at similar answers [39]. The calculation of Appelquist and Triantaphyllou, via Schwinger-Dyson equations for the technicondensates and technifermion masses, yields:

$$S \simeq (0.12)N_{TC}N_{TD}(1 - N_{TD}/8) + 0.13, \quad (12)$$

for  $N_{TD} \leq 6$ , not very different from the QCD-like result. The elegant calculation of Sundrum and Hsu makes use of the operator product expansion and less restrictive assumptions about the spectrum; they obtain:

$$S \simeq (0.11a - 0.07b)N_{TC}N_{TD} + 0.13, \quad (13)$$

where  $a, b$  are unknown constants of order unity. Again, the result is not very different from the QCD-like case. The conventional technicolor theories probably always yield positive contributions to  $S$ , at least from the resonances, because the resonant part of  $S$  varies as the difference of the axial-vector and vector resonance mass-squares; the former seems to be always larger than the latter by general properties of symmetry-breaking. It thus seems reasonable to conclude that the conventional technicolor models are disfavored by the electroweak data.

But these models are not the only possibility for dynamical symmetry breaking. Apart from searching for theories with the right mass hierarchies of pseudoGoldstone scalars [34, 40], one can drop the assumption that the technicolor force is unbroken and confining, or is vector-like, or acts

only on the exotic technifermions. Recent works [41] have introduced technicolor sectors that couple to part or all of the third generation of ordinary fermions ( $\nu_\tau, \tau, t, b$ ). This idea is appealing because it naturally explains why the top quark mass is so large compared to the masses of the other fermions. Another possibility, which can be combined with “topcolor”, is that the technicolor gauge symmetry is broken, so that it binds but does not confine and produces scalar but not vector states [42]. The low-energy interactions among technifermions would be via Fermi-like four-fermion couplings. One way for technicolor to be broken is for the technifermion content to be chiral, rather than vector-like, so that the gauge symmetry breaks itself. Any combination of these ideas may be capable of producing a minimal technicolor theory that is compact and natural, with small contributions to  $S$ .

One of the many virtues of supersymmetry (SUSY) is that it has a minimal, standard version, the so-called minimal supergravity (SUGRA) model [36]. The supersymmetry must be broken, of course, so that particles and their respective superpartners do not have the same masses. In the minimal case, this is accomplished by soft breaking terms of dimension three that endow the scalar quarks and leptons (squarks and sleptons) with a large common mass (the gravitino mass) that splits them from the fermions, and the gauge fermions (gauginos: winos, sinos, photinos, and gluinos) with their own large Majorana masses. The superpartner masses also receive contributions from electroweak symmetry breaking, just like the ordinary particles. Because the breaking of supersymmetry and the breaking of the electroweak symmetry have nothing to do with each other (at least directly), the supersymmetry-breaking contributions to the superpartner masses must be  $SU(2)_L \times U(1)_Y$ -invariant. That being the case, if the superpartner masses are larger than those of the ordinary particles, they must be dominated by the supersymmetry-breaking but gauge-invariant contributions; therefore, in the limit of heavy superpartners, *the radiative effects of supersymmetry in  $S, T, U$  naturally decouple*. Thus supersymmetry is consistent with the electroweak data in a negative way, merely by the absence of large deviations of the  $S, T, U$  from zero. (The negative value for  $S$  still poses a puzzle, but is consistent with zero at two standard deviations.) In Table 2, I show a supersymmetry-versus-technicolor checklist, with the current score distinctly, but not decisively, in favor of supersymmetry. Readers interested in the supersymmetric loop contributions to the  $S, T, U$  should consult the work of Lynn, and Barbieri *et al.* [43]. The minimal supersymmetric Standard Model must have two Higgs doublets, with one light neutral Higgs of mass less than about 150 GeV [44]; the precision fit then yields for the top quark mass:  $m_t = 145 \pm 24$  GeV [26].

	<i>Supersymmetry</i>	<i>Technicolor</i>
<i>Theoretical Principles</i>		
Gauge hierarchy	✓	✓
Minimal theory	✓ (SUGRA)	⊗
GUTs	✓	⊗
Gravity	✓✓	
<i>Radiative Limits</i>		
Electroweak	✓	⊗ ?
FCNCs	✓	✓ (WTC)
GUTs	✓✓ (SUSY $SU(5)$ )	⊗
<i>Direct Production</i>		
New states	nothing yet	nothing yet

Table 2. A comparative scorecard: supersymmetry versus technicolor.

The technicolor-versus-supersymmetry rivalry may be understood as a competition between two different paradigms for the final state of elementary particle physics, the “bang” versus the

“whimper” [45]. In the “whimper” scenario, many or all elementary particles are composite, and the constituents themselves are perhaps in turn composites...elephants all the way down. The binding or confinement of the constituents is essentially non-perturbative. Accelerators may be likened to onion-peelers in world where no onion layer is the final one. From a theoretical point of view, this scenario (which is reminiscent of the strong-interaction “bootstrap” of the 1960’s [46]) is problematic, because it is impossible to specify the degrees of freedom, the laws of motion, and the fundamental constants of the theory in closed form. Whether or not the theory is self-inconsistent or tautological becomes impossible to decide. The “bang” scenario posits that there are a finite number of irreducible degrees of freedom, specified most naturally in terms of symmetry groups, a finite number of elementary constants, and laws of motion with a closed form. Supersymmetry fits nicely into this picture, as it automatically includes gravity (when the supersymmetry is gauged), leads naturally to superstrings (the only candidate so far for a consistent theory of everything), is easy to combine with grand unification, and leads to quasi-perturbative renormalisation group structures for the gauge couplings [29] and fermion masses and mixings [47], an approach successful so far.

The fact that the  $S, T, U$  are close to, and consistent with, zero means that, whatever non-Standard physics lies beyond the  $Z$ , it must satisfy at least one of three general scenarios. The first is that the new physics is minimal to non-existent; for example, if the Higgs sector consists only of deeply bound top quark pairs. The second is that the new physics respects the  $SU(2)_L \times U(1)_Y$  gauge symmetry, with supersymmetry as a natural example. The third is that the new physics respects the larger global  $SU(2)_L \times SU(2)_R$  weak chiral symmetry (apart from the static Higgs vacuum), ensuring automatic decoupling. Implementing this symmetry naturally may be the only way to save technicolor. And non-Standard physics must also satisfy the constraints imposed by limits on flavor-changing neutral currents [21] and the requirement that the Higgs sector be radiatively stable.

## 5. FINAL THOUGHTS AND FUTURE PROSPECTS

To summarise: The presence of non-decoupled radiative effects from heavy states with weak isospin is made possible in the Standard Model by the breaking of global and gauge symmetries and can be analysed in a simple, general way through the symmetry properties of electroweak currents. The current data, all from four-fermion experiments, are now good enough to place highly non-trivial constraints on new physics beyond the  $Z$ : either there is little or no new physics, or the new physics is restricted by global and/or gauge symmetries.<sup>5</sup>

Four-fermion experiments remain that can improve our knowledge of the universal corrections. In particular, the  $W$  mass measurement at the Tevatron [49] and LEP200 [50], the polarised electron beam  $Z$  asymmetry at the SLC [51], a new atomic parity violation experiment [52], and more precise low-energy neutrino scattering measurements [53] seem the most promising areas for the final refinement of electroweak gauge boson properties, reducing the uncertainties of Figure 1. The current precision of data makes radical change in the measured  $S, T, U$  unlikely at this time. The measurement of the top quark mass, while not by itself changing the uncertainties in the precision electroweak measurements, will shift the origin of the  $S, T, U$  space and allow a clean test of the minimal Standard Model.

Particle physics until now has been a world of four-fermion interactions, mediated by gauge boson exchanges as the manifestations of elementary fields of force. The current generation of high-energy colliders — LEP (CERN), SLC (SLAC), Tevatron (Fermilab), and HERA (DESY) — will close out the era of high-energy four-fermion processes with precision electroweak measurements and

<sup>5</sup>A caveat: It is possible to add new gauge interactions, such as an extra  $Z$ . This possibility must be analysed in a framework more general than the one presented here, with additional parameters, because not all of the effects of new gauge bosons at low energies can be absorbed into the universal  $SU(2)_L \times U(1)_Y$  parameters  $S, T, U$ . The lower mass bounds on  $Z$ ’s and  $W$ ’s are now quite strong [48].

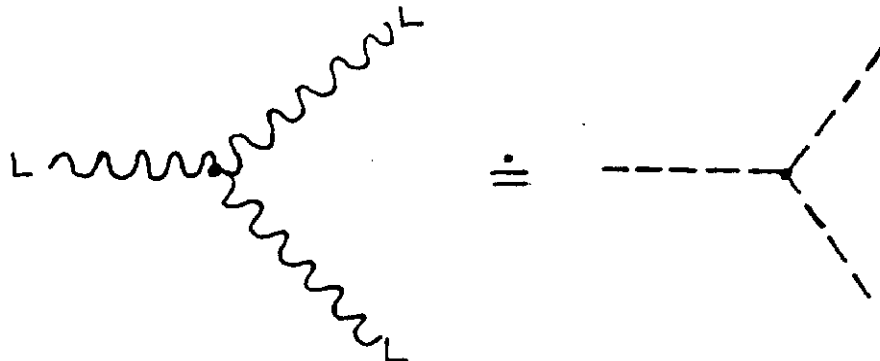


Figure 4. Electroweak gauge boson self-interactions.

probably the discovery of the top quark. The next generation of colliders — LEP200 (CERN) [50], LHC (CERN) [54], SSC [55], and possibly a new  $e^+e^-$  machine [56] — will open a new era devoted to the study of the Higgs sector and the non-Abelian electroweak gauge interactions. These new machines will give us access to interactions never directly studied before: non-Abelian self-interactions of the  $W$  and  $Z$  gauge bosons, interactions of the Higgs sector with itself and with the gauge and fermion sectors, and therefore insight into the origin of mass [57].

Since the longitudinal gauge boson modes are really Goldstone modes, the study of their self-interactions will already give us a window into the Higgs self-interactions (Figure 4). The radiative corrections to these couplings give us a new class of direct or non-Born-like corrections and thus possibly new effects arising from non-Standard sectors. The universal corrections to gauge exchange are already known from four-fermion measurements, as presented above. One four-fermion process, the production of  $b\bar{b}$  (bottom) pairs from the decay of the  $Z$ , contains a special non-universal correction from the presence of virtual top quark states (Figure 5). The relevant gauge modes in the vertices are longitudinal = disguised Goldstones again. This special effect allows us to use the  $Z \rightarrow b\bar{b}$  data to place a  $\rho$ - or  $T$ -independent bound on the top quark mass:  $m_t < 320$  (350) GeV at 90 (95)% C.L., with an independent value of  $\rho = 1.007 \pm 0.011$  [26]. The custodial  $SU(2)_V$  global subgroup seems safe.

Although the  $SU(2)_L \times U(1)_Y$  gauge symmetry and the  $SU(2)_L \times SU(2)_R$  global symmetry have been established as valid at low energies to a few tenths of a percent, important questions about the electroweak interactions remain. The most crucial of these questions concern the nature of the Higgs sector — what it is, where it comes from, and why the electroweak gauge symmetry breaks at a scale so much smaller than the grand unification and Planck scales. Related questions include: Why do the fermions have the masses and mixings that they do? Why is the top quark mass so much larger than the other fermion masses? Why is the  $SU(2)_L \times SU(2)_R$  global symmetry such a good symmetry and where does it come from? Are there any extended gauge groups, new weak gauge bosons at masses higher than the  $W$  and  $Z$ ? The precision measurement of electroweak gauge interactions, like the study of flavor-changing neutral currents and CP violation, has placed stringent limits on what non-Standard physics may lie beyond the  $Z$ . But the discovery of that physics still requires the production of the new states and new accelerators.

## ACKNOWLEDGMENTS

For use of their data analyses, I wish to thank Paul Langacker and Nir Polonsky of the University of Pennsylvania and Ming Xing Luo of the University of Washington. Steven Errede of the University

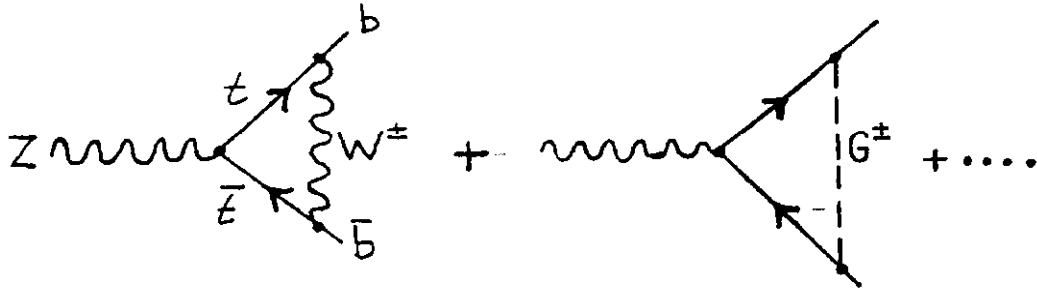


Figure 5. Non-universal top quark mass dependence in process  $Z \rightarrow b\bar{b}$ .

of Illinois (CDF) provided helpful information about the top quark. I thank also Behram Kursunoglu and Sydney Meshkov of the Global Foundation for the opportunity to present this talk at the Coral Gables Conference and all my previous audiences of the last two and a half years for listening to earlier versions. This work was supported by the U.S. Department of Energy under contract no. DOE-AC02-76-CHO-3000.

## REFERENCES

1. D. C. Kennedy, In: *Proceedings of the 1991 Theoretical Advanced Studies Institute in Elementary Particle Physics* (R. K. Ellis, C. T. Hill, and J. D. Lykken, Eds.), Singapore: World Scientific, 1992, pp. 163-280.
2. S. L. Glashow, *Nucl. Phys.*, Vol. 22, 1961, pp. 579-588; S. Weinberg, *Phys. Rev. Lett.*, Vol. 19, 1967, pp. 1264-1266; A. Salam, In: *Elementary Particle Physics* (N. Svartholm, Ed.), Stockholm: Almqvist and Wicksell, 1968, pp. 367-377; S. L. Glashow, S. Weinberg, and A. Salam, *Rev. Mod. Phys.*, Vol. 52, 1980, pp. 515-543.
3. *Quantum Electrodynamics* (J. Schwinger, Ed.), New York: Dover Publications, 1958; *Quantum Electrodynamics* (T. Kinoshita, Ed.), Singapore: World Scientific, 1990.
4. T.-P. Cheng and L.-F. Li, *Gauge Theory of Elementary Particle Physics*, New York: Oxford University Press, 1984.
5. P. Higgs, *Phys. Rev.*, Vol. 145, 1966, pp. 1156-1163; T. W. B. Kibble, *Phys. Rev.*, Vol. 155, 1967, pp. 1554-1561; J. F. Gunion, H. E. Haber, G. Kane, and S. Dawson, *The Higgs Hunter's Guide*, Redwood City, CA: Addison-Wesley, 1990.
6. *The Standard Model Higgs Boson* (M. B. Einhorn, Ed.), Amsterdam: North-Holland, 1991.
7. D. Schaile and P. M. Zerwas, *Phys. Rev.*, Vol. D45, 1992, pp. 3262-3265.
8. F. Abe *et al.* (CDF), *Phys. Rev.*, Vol. D45, 1992, pp. 3941-3948.
9. *Fermilab Program Through 1997 and Beyond: Supplemental Materials submitted to the 1992 HEPAP Subpanel*, Batavia, IL: Fermilab, 1992, pp. 3.1-3.14.
10. B. W. Lynn, M. E. Peskin, and R. G. Stuart, In: *Physics at LEP*, CERN 86-02 (J. Ellis and R. Peccei, Eds.), Geneva: CERN, 1986, Vol. 1, pp. 90-152.
11. D. C. Kennedy and B. W. Lynn, *Nucl. Phys.*, Vol. B322, 1989, pp. 1-54.

12. G. Degrandi and A. Sirlin, *Nucl. Phys.*, Vol. B383, 1992, pp. 73-92; *Phys. Rev.*, Vol. D46, 1992, pp. 3104-3116.
13. T. Appelquist and J. Carassone, *Phys. Rev.*, Vol. D11, 1975, pp. 2856-2861; J. C. Collins, *Renormalization*, Cambridge: Cambridge University Press, 1984.
14. M. Veitman, *Nucl. Phys.*, Vol. B123, 1977, pp.89-99.
15. J. C. Collins, F. Wilczek, and A. Zee, *Phys. Rev.*, Vol. D18, 1978, pp. 242-247.
16. S. Bertolini and A. Sirlin, *Nucl. Phys.*, Vol. B248, 1984, pp. 589-614.
17. W. Marciano and A. Sirlin, *Phys. Rev.*, Vol. D22, 1980, pp. 2695-2717; *Nucl. Phys.*, Vol. B189, 1981, pp. 442-460.
18. A. Sirlin, *Phys. Lett.*, Vol. 232B, 1989, pp. 123-126; S. Fanchiotti and A. Sirlin, *Phys. Rev.*, Vol. D41, 1990, pp. 319-322; A. Sirlin, *Nucl. Phys.*, Vol. B332, 1990, pp. 20-38 G. Degrandi, S. Fanchiotti, and A. Sirlin, *Nucl. Phys.*, Vol. B351, 1991, pp. 49-69.
19. D. C. Kennedy, *Phys. Lett.*, Vol. 268B, 1991, pp. 86-88.
20. T. Appelquist and C. Bernard, *Phys. Rev.*, Vol. D22, 1980, pp. 200-213; A. C. Longhitano, *Phys. Rev.*, Vol. D22, 1980, pp. 1166-1175; *Nucl. Phys.*, Vol. B188, 1981, pp. 118-154.
21. S. L. Glashow, J. Iliopoulos, and L. Maiani, *Phys. Rev.*, Vol. D2, 1970, pp. 1285-1292; F. J. Gilman and Y. Nir, *Ann. Rev. Nucl. Part. Sci.*, Vol. 40, 1990, pp. 213-238; J. L. Rosner, In: *Proceedings of the 1990 Theoretical Advanced Studies Institute in Elementary Particle Physics* (M. Cvetič and P. Langacker, Eds.), Singapore: World Scientific, 1991, pp. 91-224.
22. M. E. Peskin and T. Takeuchi, *Phys. Rev. Lett.*, Vol. 65, 1990, pp. 964-967; *Phys. Rev.*, Vol. D46, 1992, pp. 381-409.
23. D. C. Kennedy and P. Langacker, *Phys. Rev. Lett.*, Vol. 65, 1990, pp. 2967-2970; *Phys. Rev.*, Vol. D44, 1991, pp. 1591-1592; W. Marciano and J. L. Rosner, *Phys. Rev. Lett.*, Vol. 65, 1990, pp. 2963-2966; G. Altarelli and R. Barbieri, *Phys. Lett.*, Vol. 253B, 1991, pp. 161-167.
24. U. Amaldi *et al.*, *Phys. Rev.*, Vol. D36, 1987, pp. 1385-1407.
25. M. C. Noecker, B. P. Johnson, and C. E. Weiman, *Phys. Rev. Lett.*, Vol. 61, 1988, pp. 310-313; E. R. Boston and P. G. H. Sandars, *J. Phys.*, Vol. B23, 1990, pp. 2663-2672; S. A. Blundell, J. Sapirstein, and W. R. Johnson, *Phys. Rev.*, Vol. D45, 1992, pp. 1602-1623.
26. P. Langacker, U. Pennsylvania preprint UPR-0492-T, 1992, 42pp., and private communication; S. Weinberg, U. Texas preprint UTTG-25-92, 1992, 39pp.
27. A. Djouadi, *Nuovo Cim.*, Vol. 100A, 1988, pp. 357-371.
28. P. Langacker and M. X. Luo, *Phys. Rev.*, Vol. D44, 1991, pp. 817-822.
29. U. Amaldi, W. de Boer, and H. Furstenau, *Phys. Lett.*, Vol. B260, 1991, pp. 447-455; P. Langacker and N. Polonsky, U. Pennsylvania preprint UPR-0513-T, 1992, 57pp.; R. K. Ellis, Fermilab preprint FERMILAB-CONF-93/011-T, 1993, 17pp.
30. D. C. Kennedy and B. W. Lynn, SLAC preprint SLAC-PUB-4608, 1988, 18pp.
31. Y. Totsuka, In: *Proceedings, Underground Physics Experiments* (K. Nakamura, Ed.), Tokyo: U. Tokyo, 1990, pp. 129-163.
32. The CDF and D0 detectors, by the end of the third Tevatron run, will be able to place a lower limit of 250 (180) GeV on the top quark mass with (without) the Main Injector.
33. E. Gates and J. Terning, *Phys. Rev. Lett.*, Vol. 67, 1991, pp. 1840-1843; S. Bertolini and A. Sirlin, *Phys. Lett.*, Vol. 257B, 1991, pp. 179-186.
34. M. Dugan and L. Randall, *Phys. Lett.*, Vol. 264B, 1991, pp. 154-160.
35. S. Weinberg, *Phys. Rev.*, Vol. D19, 1979, pp. 1277-1280; L. Susskind, *Phys. Rev.*, Vol. D20, 1979, pp. 2619-2625; S. Dimopoulos and L. Susskind, *Nucl. Phys.*, Vol. B155, 1979, pp.

- 237-252; E. Eichten and K. Lane, *Phys. Lett.*, Vol. 90B, 1980, pp. 125-130; R. K. Kaul, *Rev. Mod. Phys.*, Vol. 55, 1983, pp. 449-475.
36. H. P. Nilles, *Phys. Rep.*, Vol. 110C, 1984, pp. 1-162; H. E. Haber and G. Kane, *Phys. Rep.*, Vol. 117C, 1985, pp. 76-263; R. N. Mohapatra, *Unification and Supersymmetry*, New York: Springer-Verlag, 1986.
  37. B. Holdom, *Phys. Lett.*, Vol. 150B, 1985, pp. 301-305; Vol. 198B, 1987, pp. 535-542; T. Appelquist, D. Karabali, and L. C. R. Wijewardhana, *Phys. Rev. Lett.*, Vol. 57, 1986, pp. 957-960.
  38. M. Golden and L. Randall, *Nucl. Phys.*, Vol. B361, 1991, pp. 3-23.
  39. R. Sundrum and S. D. H. Hsu, U. California at Berkeley preprint UCB-PTH-91-34, 1991, revised 1992, 35pp.; T. Appelquist and G. Triantaphyllou, *Phys. Lett.*, Vol. 278B, 1992, pp. 345-350.
  40. M. A. Luty and R. Sundrum, Lawrence Berkeley preprint LBL-32893-REV, 1992, 10pp.
  41. W. A. Bardeen, C. T. Hill, and M. Lindner, *Phys. Rev.*, Vol. D41, 1990, pp. 1647-1660; W. J. Marciano, *Phys. Rev.*, Vol. D41, 1990, pp. 219-224; C. T. Hill, *Phys. Lett.*, Vol. 266B, 1991, pp. 419-424; S. P. Martin, *Phys. Rev.*, Vol. D45, 1992, pp. 4283-4293; *Phys. Rev.*, Vol. D46, 1992, pp. 2197-2202.
  42. C. T. Hill, M. Luty, and E. A. Paschos, *Phys. Rev.*, Vol. D43, 1991, pp. 3011-3025; M. A. Luty, Lawrence Berkeley preprint LBL-32089, 1992, 24pp.; R. Sundrum, Lawrence Berkeley preprint LBL-32107, 1992, 21pp.; C. T. Hill, D. C. Kennedy, T. Onogi, and H.-L. Yu, Fermilab preprint FERMILAB-PUB-92/218-T, 1992, 23pp.; S. P. Martin, Northeastern U. preprint NUB-3058-92TH, 1992, 18pp.
  43. B. W. Lynn, SLAC preprint SLAC-PUB-3358, 1984, 40pp.; R. Barbieri *et al.*, *Nucl. Phys.*, Vol. B341, 1990, pp. 309-321.
  44. H. Haber and R. Hempfling, *Phys. Rev. Lett.*, Vol. 66, 1991, pp. 1815-1818.
  45. I owe this formulation to Paul Langacker.
  46. G. Chew, *The Analytic S-Matrix*, New York: W. A. Benjamin, 1964.
  47. H. Arason *et al.*, *Phys. Rev.*, Vol. D46, 1992, pp. 3945-3965; *Phys. Rev.*, Vol. D47, 1992, pp. 232-240; S. Dimopoulos, L. J. Hall, and S. Raby, *Phys. Rev.*, Vol. D46, 1992, pp. 4793-4798; U. California at Berkeley preprint UCB-PTH-92-21, 1992, 16pp.; Ohio State U. preprint OHSTPY-HEP-92-018, 1992, 25pp.
  48. P. Langacker and M. X. Luo, *Phys. Rev.*, Vol. D45, 1992, pp. 278-292; F. del Aguila *et al.*, *Nucl. Phys.*, Vol. B361, 1991, pp. 45-71; *Nucl. Phys.*, Vol. B372, 1992, pp. 3-22.
  49. The combined error on the  $W$  mass from CDF and D0 will be 35 (70) MeV with (without) the Main Injector. See ref. [9].
  50. The LEP200 error on the  $W$  mass will be 100 MeV. See: G. Barbiellini *et al.*, In: *Physics at LEP*, CERN 86-02 (J. Ellis and R. Peccei, Eds.), Geneva: CERN, 1986, Vol. 2, pp. 1-79.
  51. B. W. Lynn and R. G. Stuart, *Nucl. Phys.*, Vol. B253, 1985, pp. 216-230; SLC Polarisation Group, *SLC Polarization Proposal*, SLAC report SLAC-PROPOSAL-SLC-01, 1986, 176pp.; B. W. Lynn and C. Versegnassi, *Phys. Rev.*, Vol. D35, 1987, pp. 3328-3337. With 50,000  $Z$ 's at 60% polarisation,  $\sin^2 \theta_W$  can be measured to  $\pm 0.0008$ .
  52. C. E. Weiman, In: *Proceedings of the 1989 Theoretical Advanced Studies Institute in Elementary Particle Physics* (T. DeGrand and D. Toussaint, Eds.), Singapore: World Scientific, 1990, pp. 645-654; J. Sapirstein, *ibid.*, pp. 655-674. An improvement of the error on  $S$  from atomic measurements from  $\pm 4$  to  $\pm 1.5$  is not unrealistic.
  53. T. Bolton *et al.*, Fermilab proposal P-815, 1990, 92pp. The P-815 collaboration envisions



measuring  $\sin^2 \theta_W$  to  $\pm 0.0025$ .

- 54. *Design Study of the Large Hadron Collider (LHC)*, CERN 91-03, Geneva: CERN, 1991.
- 55. *Research Directions for the Decade: Snowmass '90 Proceedings* (E. L. Berger, Ed.), Singapore: World Scientific, 1992.
- 56. C. Ahn *et al.*, *Opportunities and Requirements for Experimentation at High Energy  $e^+e^-$  Collider*, SLAC report SLAC-329, 1988, 224pp.
- 57. J. A. Bagger, In: *Proceeding of the 1991 Theoretical Advanced Studies Institute in Elementary Particle Physics* (R. K. Ellis, C. T. Hill, and J. D. Lykken, Eds.), Singapore: World Scientific, 1992, pp. 31-94; C. P. Burgess and D. London, *Phys. Rev. Lett.*, Vol. 69, 1992, pp. 3428-3431.